

Transport Catastrophe Near the Superconductor-Insulator Transition

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Thin films of Amorphous indium oxide undergo a magnetic field driven superconducting to insulator quantum phase transition. In the insulating phase, the current-voltage characteristics show large current discontinuities due to overheating of electrons. Here we follow the temperature and magnetic field evolution of the discontinuities while approaching the quantum critical point. We show that the onset voltage for the discontinuities vanishes as we approach the quantum critical point. As a result the insulating phase becomes unstable with respect to any applied voltage making it, at least experimentally, immeasurable. We emphasize that unlike previous reports of the absence of linear response near quantum phase transitions, in our system, the departure from equilibrium is discontinuous. We discuss the "catastrophic" implications of such a discontinuous response on transport measurements. Whether the instability is only an experimental barrier or also a theoretical limit remains an open question.

The superconducting to insulator transition (SIT) [1, 2], observed in thin films of highly disordered superconductors, is a quantum phase transition (QPT) [3] where by varying the magnetic field (B) [4–6], level of disorder [7], film thickness [8] or charge carrier density [9], the ground state of the system transforms from superconducting to insulating.

In the B -driven SIT, beyond the critical B (B_c), Cooper-pairs persist and become spatially localized [10–16], leading to the emergence of a strongly insulating state [5, 12, 14, 17–19]. Beyond B_c , the resistance (R) sharply increases with B , then reaches a maximum at a T -dependent B , B_{peak} , above which R begins to drop. Recently we reported that at $B_c < B < B_{peak}$, R appears to diverge at a finite T [20].

In the insulating phase, at low T (typically $T \leq 200$ mK), the current-voltage characteristics ($I-V$'s) exhibit large I -discontinuities (ΔI) [13] (see figure 1). The $I-V$'s can be separated into two distinct regions, a high resistance (HR) state at low V and a low resistance (LR) state at high V . The transition between the two states occurs at two well-defined V 's. Increasing V from $V = 0$ results in a ΔI at a T dependent V , V_{escape} , where the sample is driven from the HR state to the LR state. Decreasing V from the LR state result in another ΔI at a T independent V , V_{trap} , where the $I-V$'s switch from the LR to the HR state. This process is hysteretic and $V_{trap} \leq V_{escape}$.

These ΔI 's were first associated with the emergence of a new "superinsulating" phase dual to superconductivity [21]. Later Altshuler *et al.* took a different approach and showed that the ΔI 's could be a manifestation of a thermal bi-stability where the electrons abruptly attain a T much greater than the phonon T (T_{ph}) [22]. The central assumption of this model is that all deviations from a linear $I-V$ result from an increase in electron T (T_{el}). T_{el} is determined by the heat-balance between the

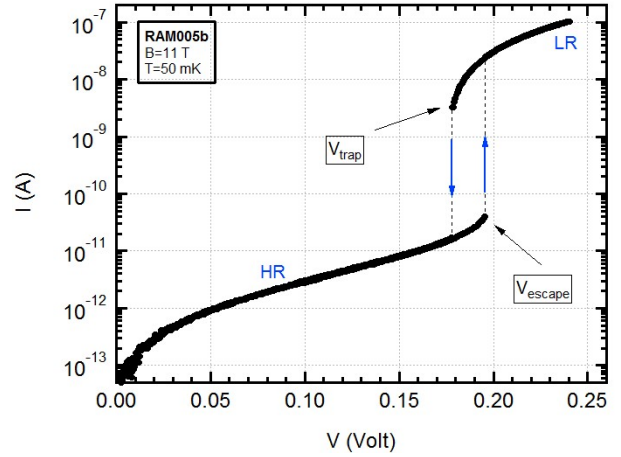


FIG. 1: **Discontinuities in the $I-V$'s.** I (log scale) vs. V measured at $B = 11$ T (in the insulating phase) and $T = 50$ mK. The measured data points are marked by full circles, the dashed line connecting the data points is a guide for the eye. The LR→HR (HR→LR) transition is marked by a blue arrow pointing down (up). At $V = 0$ the sample is in the HR state. By increasing V the $I-V$'s exhibit a discontinuity at $V_{escape} = 0.195$ Volts where I jumps by 3 orders of magnitude. Decreasing V results in a hysteresis where I drops back to the HR state at $V_{trap} = 0.178$ V.

experimentally applied Joule-heating ($I \cdot V$) and cooling via the phonons. To allow for T_{el} that is different from phonon T (T_{ph}) they assumed that the electron-phonon coupling vanishes rapidly at low T 's (as observed in various materials [23, 24]). This non-equilibrium state is analyzed, under steady state conditions, by solving the

heat-balance equation

$$\frac{V^2}{R(T_{el})} = \Gamma \Omega (T_{el}^\beta - T_{ph}^\beta) \quad (1)$$

where Ω is the volume of the sample, Γ is the electron-phonon coupling-strength, β is an exponent that determines the power-law decay of the electron-phonon coupling as $T \rightarrow 0$ (taken to be 6, which is appropriate for a metal in the dirty limit [22, 25, 26]) and $R(T_{el}) = R_0 \exp((T_0/T_{el})^\gamma)$, typical for insulators.

The central result is that, below a critical T_{ph} , the heat-balance equation has two stable solutions for T_{el} ; The ΔI 's are a result of a change in R that occurs when the electrons abruptly switch between the low T_{el} solution, where a HR state exists, and the high T_{el} solution, resulting in a far-from equilibrium state for the system. This electron-heating approach gained support from several experiments [27–29].

At first sight, the ΔI 's and the associated electron-phonon decoupling is not relevant to the study of the underlying equilibrium phases, and of the SIT itself. Upon deeper inspection, it appears that this is not the case. In this letter we systematically followed the B -evolution of the ΔI and found that V_{escape} vanishes as $B \rightarrow B_c$. Consequently, close enough to B_c , the finite V required for transport measurements will inevitably exceed V_{escape} , driving the system out of equilibrium with a discontinuous transition to the high T_{el} state. The significance of the discontinuous departure from equilibrium is that, not only the equilibrium state cannot be measured but, near the SIT, it is not experimentally possible to extrapolate the equilibrium properties from the measurable R .

Our data were obtained by measuring two thin films of amorphous indium-oxide. The films were deposited by e-gun evaporation of high purity In_2O_3 onto a SiO_2 substrate in an O_2 rich environment. Both samples have a Hall-bar geometry, their lengths and widths are $2 \times 0.5 \text{ mm}^2$ (sample RAM005b) and $10 \times 5 \mu\text{m}^2$ (sample BT1cH5) and their thickness is 30 nm. The data presented was measured in a two-probe dc configuration.

The central results of this work are presented in figure 2, where we display V_{escape} vs. $\delta B \equiv \frac{B-B_c}{B_c}$ at $T = 11 \text{ mK}$. On this log-log plot V_{escape} follows a power-law that spans almost two decades in δB and four decades in V_{escape} , indicating that V_{escape} vanishes rapidly upon approaching B_c .

$$V_{escape}(B) \propto (B - B_c)^\alpha \quad (2)$$

where $\alpha = 2.24$ is extracted using a power-law fit (α appears to be non-universal: it is sample and T -dependent.).

A vanishingly small V_{escape} upon approaching B_c imposes strong constraints for a theoretical description of our disordered films. Some theories describe similar discontinuous $I - V$'s and their V threshold [21, 30, 31] in

terms of Coulomb blockade on an assembly of grains, for which V_{escape} would relate to the sum of the charging energies of each grain. However, V_{escape} as low as $10 \mu\text{V}$ at $\delta B \sim 0.1$ in fig 2 would results in a vanishingly small charging energy, for any reasonable assumption of the number of grains. Therefore V_{escape} in our system cannot comply with such a Coulomb blockade picture and rules out models based on granular structures.

The central implication of equation 2 is that it reflects an inherent experimental difficulty one is faced when measuring equilibrium properties near B_c . While conducting transport measurements, it is essential to apply a finite V across the sample. This applied V must exceed the noise level present during the experiment (such as Johnson-Nyquist noise) and must also be large enough to induce a measurable I response from the sample (typically $V > \mu\text{V}$). The vanishing of V_{escape} suggests a loosing cause: whatever small V is, there will always be a B range, close to the SIT, where it will exceed V_{escape} and drive the system out of equilibrium. It is known that systems exhibit a non-linear response near QPT's [32–34] where, at $T = 0$, any finite V will drive them out of equilibrium. The pivotal difference reported in this letter is that, not only our system has no linear response, but it departs from equilibrium in a discontinuous fashion. In the discussion section we present several consequences of this discontinuous response.

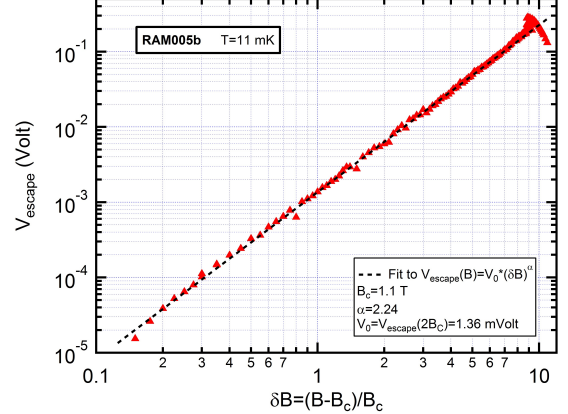


FIG. 2: **Magnetic-field evolution of V_{escape} .** V_{escape} vs $\delta B \equiv \frac{B-B_c}{B_c}$ (log-log scale) at $T = 11 \text{ mK}$. The data are presented as red triangles while a dashed black-line marks a power-law fit.

Because, strictly speaking, the QPT takes place at $T = 0$ it is worthwhile to examine the T evolution of V_{escape} . If V_{escape} increases as $T \rightarrow 0$, at low enough T 's, the HR state might span a large enough V interval and become measurable. In Figure 3a we display the $I - V$'s measured at $B = 9.5 \text{ T}$ at different T 's. Below $T = 100 \text{ mK}$, the $I - V$'s become discontinuous. As predicted [22], V_{trap} is nearly independent of T_{ph} . V_{escape} ,

on the other hand, initially increases as T_{ph} is reduced down to $T_{ph} = 60$ mK as expected. As T_{ph} is further lowered, the trend changes and V_{escape} begins to decrease. In figure 3b we display V_{escape} and V_{trap} (up and down pointing triangles respectively) vs T . It appears that $\lim_{T_{ph} \rightarrow 0} V_{escape}(T_{ph}) \sim V_{trap}$. In the inset of figure 3b we display V_{escape} and V_{trap} vs T at various B 's. For all measured B 's, V_{escape} follows a similar pattern. These results suggest that lowering T will not increase V_{escape} and make the HR state measurable. In the discussion we show that this T dependence of V_{escape} can be explained by considering the inhomogeneity of the system.

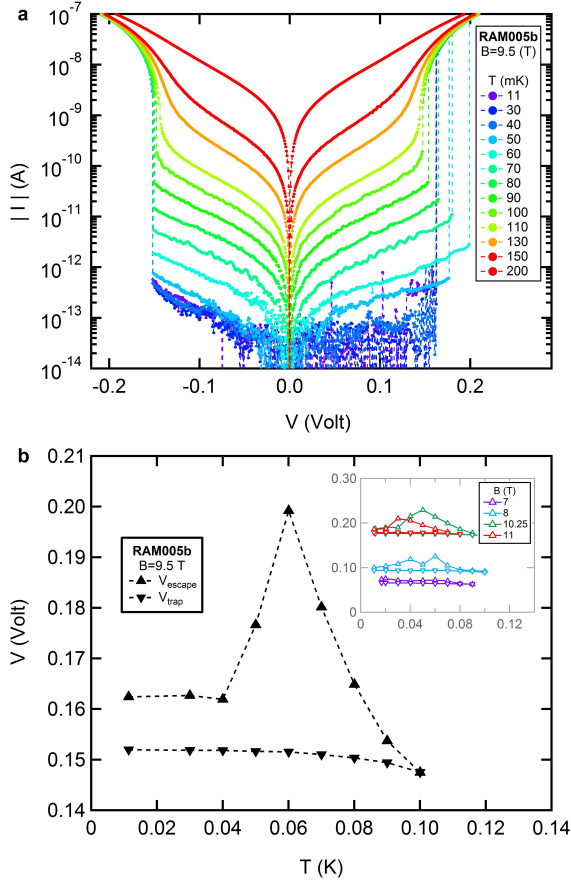


FIG. 3: **Saturation of V_{escape} as $T_{ph} \rightarrow 0$.** (a) $|I|$ (log scale) vs. V measured at $B = 9.5$ T. The color coding describe different T_{ph} isotherms ranging from 11 mK (purple) to 200 mK (red). At $T_{ph} = 100$ mK, the $I - V$'s become discontinuous with $V_{escape} \sim 150$ mV. While cooling, V_{escape} initially increases up to $V_{escape} = 200$ mV at $T_{ph} = 60$ mK. At lower T_{ph} , V_{escape} drops and saturates at a finite V . (b) V_{escape} and V_{trap} vs. T_{ph} at $B = 9.5$ T. V_{escape} and V_{trap} are marked by up and down pointing triangles respectively. At low T 's, V_{escape} saturates at a value which is comparable to V_{trap} . Inset: V_{escape} and V_{trap} vs. T_{ph} at different B 's.

An important question is how does the magnitude of

the ΔI 's evolve while approaching the quantum critical point (QCP). If the ΔI 's vanishes sufficiently fast, the transition from the LR to the HR states becomes practically continuous in the sense that one might still be able to extrapolate equilibrium, HR, properties from the LR state. The B -dependence of the ΔI 's at $T = 11$ mK is displayed in figure 4a where we focus on the trapping side where the LR \rightarrow HR transition occurs. The blue triangles correspond to the left (blue) axis and represent I on both sides of V_{trap} where the upwards pointing triangles correspond to the last measured I in the LR state before the jump (I_{LR}) and downwards pointing triangles stand for I_{HR} , the first measured I in the HR state (for most B 's I_{HR} was in the noise level and should only be considered as an upper bound). The red triangles mark V_{trap} and correspond to the right (red) axis. While V_{trap} vanishes rapidly over a vast B range, the magnitude of the ΔI 's does not seem to vary significantly. This observation has a great impact on the reliability of I -bias transport measurements as we will argue in the discussion. Close to B_c the magnitude of ΔI 's drops and seems to vanish, precluding a finite I flow without any applied V at B_c .

To study whether the drop in the magnitude of the ΔI 's near B_c makes it possible to extrapolate the equilibrium R from measurable LR data we define R on both sides of V_{trap} as $R_{LR} \equiv V_{trap}/I_{LR}$ and $R_{HR} \equiv V_{trap}/I_{HR}$ (we did not use dV/dI due to our inability to measure I_{HR}). In figure 4b we display R_{LR} (red) and R_{HR} (blue) vs T measured at a constant B (9.5 T). It seems that, while R_{HR} diverges as $T \rightarrow 0$, R_{LR} saturates at a finite value. In figure 4c we display R_{LR} vs δB extracted from the data of figure 4a. It can be seen that R_{LR} does not vary significantly in the vicinity of B_c and exceeds 1 M Ω only above $2B_c$. These low- R values (for an insulator at $T = 11$ mK) suggest that, approaching B_c , the LR state will not "catch up" with the diverging HR state with some sudden increase in T . Therefore, near B_c , any attempt to extrapolate $R = \lim_{V \rightarrow 0} V/I$ will yield R_{LR} that can differ by orders of magnitude from the equilibrium R .

Discussion The data presented above raises several points. Consequences relating to I biased, 4-probe, measurements - So far we discussed V biased measurements and saw that close to the SIT they will probe only the LR state and at B sufficiently far from B_c the HR state can also be probed. From figures 4a we see that in order to probe the HR state in a I biased measurement, at any $B > B_c$, one needs to apply ~ 0.1 pA where the typical I used in a I biased 4-probe measurements is 1 nA. We should note that the I values bordering the jumps are sample, T and B dependent, but typically 1 nA is either in the LR state or somewhere between the HR and LR states. In figure 5c we display data measured on sample BT1cH5, a $10 \times 5 \mu\text{m}^2$ hall bar where at $T = 15$ mK, $I = 1$ nA was in the LR state for any $B > B_c$. The red triangles

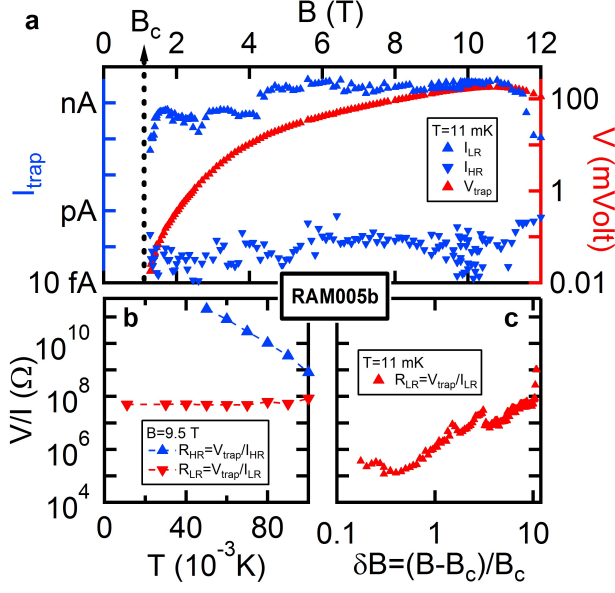


FIG. 4: T and B dependence of the ΔI 's. (a) $I_{\text{trap}} \equiv I(V_{\text{trap}})$ and V_{trap} vs. B . The blue triangles correspond to the left (blue) axis and mark I on both sides of V_{trap} where the upwards (downwards) pointing triangles correspond to the LR (HR) side of the jump. The red triangles correspond to the right (red) axis and mark V_{trap} . The vertical dashed black-line marks $B_c = 1.1$ T. (b) R_{LR} (red) and R_{HR} (blue) vs T measured at $B = 9.5$ T. (c) R_{LR} vs $\frac{B-B_c}{B_c}$ measured at $T = 11$ mK.

mark R extracted from the zero V limit of two-probe DC $I-V$'s (as previously discussed, close to B_c we probably probed only the LR state). The continuous blue line is the result of a standard AC 4-probe measurement with $I_{\text{rms}} = 1$ nA. We note that, due to the existence of the LR state, the I used in the 4-probe measurement did not drop, therefore, there was no indication that a HR state exists. The dramatic difference in R 's between the two different measurements emphasizes that, at low T 's, one should not trust measured data without first seeing the full $I-V$'s.

Another question that arises is why does V_{escape} decrease at low T 's? As displayed in figure 3, at low T 's, V_{escape} suddenly drops and saturates at a finite value which is similar to V_{trap} . Here we argue that this drop and the ensuing saturation can be accounted for by considering the effect of inhomogeneity on the electron-heating model [35]. In their theoretical work, Altshuler *et al.* noted that the ΔI 's can occur anywhere within a certain bi-stability V interval (ΔV), $V_{\text{trap}}^{\text{min}} < V < V_{\text{escape}}^{\text{max}}$, where $V_{\text{trap}}^{\text{min}}$ is a lower bound on V_{trap} and $V_{\text{escape}}^{\text{max}}$ is an upper bound on V_{escape} and they satisfy the following parametric dependences:

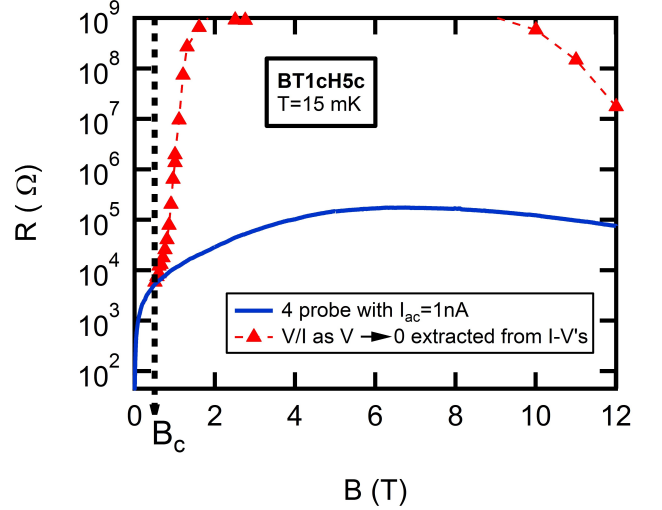


FIG. 5: Unavoidable "transport catastrophe". R (log scale) vs. B of sample BT1cH5 at $T = 15$ mK. The red triangles were extracted from the $I-V$'s in the limit $V \rightarrow 0$. The blue line was measured using a standard 4-probe AC measurement with $I_{\text{rms}} = 1$ nA.

$$V_{\text{trap}}^{\text{min}} \propto \Delta^{\frac{\beta}{2}} \quad (3)$$

$$V_{\text{escape}}^{\text{max}} \propto \Delta^{-\frac{\gamma}{2}} T_{\text{ph}}^{\frac{\beta+\gamma}{2}} e^{\frac{1}{2}(\frac{\Delta}{T_{\text{ph}}})^{\gamma}} \quad (4)$$

Because $\lim_{B \rightarrow B_c} \Delta = 0$ these equations suggest that, as $T_{\text{ph}} \rightarrow 0$ and $B \rightarrow B_c$, $V_{\text{trap}}^{\text{min}} \rightarrow 0$ and $V_{\text{escape}}^{\text{max}} \rightarrow \infty$. Therefore, for some reason, both V_{trap} and V_{escape} seem to occur prematurely i.e. at the low limit of the ΔV . One can gain intuition into why are the ΔI 's premature by considering their similarity to first order phase transitions in a Van der Waals liquid that, as we recently reported, shares several properties with our system [36]. In a Van der Waals liquid, the actual phase transition occurs not at the limit of stability but where the Maxwell area rule prescribes. Only rarely, when there are no centers of nucleation such as contamination of the fluid one can have an overcooled liquid which is close to the limit of stability.

If V_{escape} is governed by some Maxwell area law equivalent, in a manner akin to the thermodynamic case, one would expect V_{escape} to increase while cooling. This is indeed the case initially but it breaks at low T 's. The premature jumps can be explained qualitatively by assuming that some inhomogeneity exists in the sample (indications for such non-structural inhomogeneity were reported in several highly disordered superconductors [19, 37–39]). At low T 's, there is a competition between the inefficient cooling via phonons and an inefficient Joule-heating (due to the large R of the sample).

While the cooling is probably less sensitive to imperfections, the heating may be affected locally near the impurities giving rise to "hot spots" that act similarly to nucleation centers in the Van der Waals liquid. The T where V_{escape} drops should provide us with some energy scale that characterizes the sample's inhomogeneity.

Theoretical values of R_{LR} and R_{HR} while approaching the SIT - According to reference [22], at $T_{ph} = 0$, T_{el} in the LR state equals $T_{el}^{LR}(T_{ph} = 0) = \Delta(\frac{\gamma}{\beta})^{\frac{1}{\gamma}}$ while in the HR state the electrons and phonons are in thermal equilibrium therefore $T_{el}(T_{ph} = 0) = 0$. Inserting these values to the insulator's $R(T)$ (as defined earlier) we obtain that, in the HR state, R diverges while in the LR state $\lim_{T_{ph} \rightarrow 0} R^{LR}(T_{ph}) = R_0 e^{\frac{\beta}{\gamma}}$, which is finite ($\sim M\Omega$) and independent of Δ . These results are in agreement with the B -evolution of R_{LR} near the SIT and with large differences between R_{LR} and R_{HR} , as displayed in figures 4b and c.

We would like to note a prior experimental work, of Parendo *et al.* [40], that also acts on the premise that various experimental observations, near the SIT, occur due to electron heating. Their work focuses on showing that, due to electron-heating, one cannot consider the results of electric field and T scaling independently. They do not discuss whether a linear-response exists or consider a discontinuous response.

In summary, our findings questions the stability of an equilibrium, insulating, state bordering the SIT. The main result of our work is that V_{escape} , above which only the LR state persists, vanishes as we approach the QCP ($T \rightarrow 0$ and $B \rightarrow B_c$). This seemingly innocuous behavior has far-reaching consequences on transport: Because transport measurements require a finite V in order to probe the sample, only the LR state can be accessed. In addition, any V noise will already heat the electrons and drive the sample to the LR state. A central question that remains is whether, theoretically, there is a V range where the HR state is stable.

We are grateful to V. Kravtsov and I. Aleiner for fruitful discussions. This work was supported by the Israeli Science Foundation.

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- [1] A. M. Goldman and N. Markovic, *Phys. Today* **51**, 39 (1998).
- [2] V. F. Gantmakher and V. T. Dolgoplov, *Phys.-Usp.* **53**, 1 (2010).
- [3] S. L. Sondhi, S. M. Girvin, J. P. Carini, and D. Shahar, *Rev. Mod. Phys.* **69**, 315 (1997).
- [4] A. F. Hebard and M. A. Paalanen, *Phys. Rev. Lett.* **65**, 927 (1990).
- [5] A. Yazdani and A. Kapitulnik, *Phys. Rev. Lett.* **74**, 3037 (1995).
- [6] T. I. Baturina, D. R. Islamov, J. Bentner, C. Strunk, M. R. Baklanov, and A. Satta, *JETP Lett.* **79**, 337 (2004).
- [7] D. Shahar and Z. Ovadyahu, *Phys. Rev. B* **46**, 10917 (1992), URL <http://link.aps.org/doi/10.1103/PhysRevB.46.10917>.
- [8] D. B. Haviland, Y. Liu, and A. M. Goldman, *Phys. Rev. Lett.* **62**, 2180 (1989).
- [9] K. A. Parendo, K. Tan, A. Bhattacharya, M. Eblen-Zayas, N. E. Staley, and A. M. Goldman, *Phys. Rev. Lett.* **94**, 197004 (2005).
- [10] M. Feigel'man, L. Ioffe, V. Kravtsov, and E. Cuevas, *Annals of Phys.* **325**, 1390 (2010).
- [11] Y. Dubi, Y. Meir, and Y. Avishai, *Nature* **449**, 876 (2007).
- [12] V. F. Gantmakher, M. V. Golubkov, J. G. S. Lok, and A. K. Geim, *JETP* **82**, 951 (1996).
- [13] G. Sambandamurthy, L. W. Engel, A. Johansson, E. Peled, and D. Shahar, *Phys. Rev. Lett.* **94**, 017003 (2005).
- [14] H. Q. Nguyen, S. M. Hollen, M. D. Stewart, J. Shainline, A. Yin, J. M. Xu, and J. M. Valles, *Phys. Rev. Lett.* **103**, 157001 (2009).
- [15] B. Sacépé, T. Dubouchet, C. Chapelier, M. Sanque, M. Ovadia, D. Shahar, M. Feigel'man, and L. Ioffe, *Nat. Phys.* **7**, 239 (2011).
- [16] B. Sacépé, J. Seidemann, M. Ovadia, I. Tamir, D. Shahar, C. Chapelier, C. Strunk, and B. Piot, *Physical Review B* **91**, 220508 (2015).
- [17] M. A. Paalanen, A. F. Hebard, and R. R. Ruel, *Phys. Rev. Lett.* **69**, 1604 (1992).
- [18] G. Sambandamurthy, L. W. Engel, A. Johansson, and D. Shahar, *Phys. Rev. Lett.* **92**, 107005 (2004).
- [19] B. Sacépé, C. Chapelier, T. I. Baturina, V. M. Vinokur, M. R. Baklanov, and M. Sanquer, *Phys. Rev. Lett.* **101**, 157006 (2008).
- [20] M. Ovadia, D. Kalok, I. Tamir, S. Mitra, B. Sacépé, and D. Shahar, *Scientific reports* **5** (2015).
- [21] V. M. Vinokur, T. I. Baturina, M. V. Fistul, A. Y. Mironov, M. R. Baklanov, and C. Strunk, *Nature* **452**, 613 (2008).
- [22] B. L. Altshuler, V. E. Kravtsov, I. V. Lerner, and I. L. Aleiner, *Phys. Rev. Lett.* **102**, 176803 (2009).
- [23] M. L. Roukes, M. R. Freeman, R. S. Germain, R. C. Richardson, and M. B. Ketchen, *Phys. Rev. Lett.* **55**, 422 (1985), URL <http://link.aps.org/doi/10.1103/PhysRevLett.55.422>.
- [24] M. E. Gershenson, Y. B. Khavin, D. Reuter, P. Schafmeister, and A. D. Wieck, *Phys. Rev. Lett.* **85**, 1718 (2000), URL <http://link.aps.org/doi/10.1103/PhysRevLett.85.1718>.
- [25] M. Reizer and A. Sergeev, *Zh. Eksp. Teor. Fiz.* **90**, 1056 (1986).
- [26] A. Schmid, *Zeitschrift für Physik* **271**, 251 (1974).
- [27] M. Ovadia, B. Sacepe, and D. Shahar, *Phys. Rev. Lett.* **102**, 176802 (2009).
- [28] D. Kalok, A. Bilusic, T. I. Baturina, V. M. Vinokur, and C. Strunk, *arXiv* (2010), 1004.5153.v1.
- [29] T. Levinson, A. Doron, I. Tamir, C. Tawari, and D. Shahar (2016), manuscript submitted for publication.
- [30] A. A. Middleton and N. S. Wingreen, *Physical review letters* **71**, 3198 (1993).
- [31] T. I. Baturina, A. Y. Mironov, V. M. Vinokur, M. R. Baklanov, and C. Strunk, *Phys. Rev. Lett.* **99**, 257003 (2007).

- [32] A. G. Green and S. Sondhi, Physical review letters **95**, 267001 (2005).
- [33] P. Hogan and A. Green, Physical Review B **78**, 195104 (2008).
- [34] D. Dalidovich and P. Phillips, Physical review letters **93**, 027004 (2004).
- [35] V. Kravtsov, Private communication.
- [36] A. Doron, I. Tamir, S. Mitra, G. Zeltzer, M. Ovadia, and D. Shahar, Phys. Rev. Lett. **116**, 057001 (2016), URL <http://link.aps.org/doi/10.1103/PhysRevLett.116.057001>.
- [37] D. Kowal and Z. Ovadyahu, Physica C: Superconductivity **468**, 322 (2008).
- [38] D. Kowal and Z. Ovadyahu, Solid state communications **90**, 783 (1994).
- [39] L. B. Ioffe and M. E. Gershenson, Nature materials **11**, 567 (2012).
- [40] K. A. Parendo, K. S. B. Tan, and A. Goldman, Physical Review B **74**, 134517 (2006).